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Warm inflation within a supersymmetric distributed mass model

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We study the dynamics and observational predictions of warm inflation within a supersymmetric distributed mass model. This dissipative mechanism is well described by the interactions between the inflaton and a tower of chiral multiplets with a mass gap, such that different bosonic and fermionic fields become light as the inflaton scans the tower during inflation. We examine inflation for various mass distributions, analyzing in detail the dynamics and observational predictions. We show, in particular, that warm inflation can be consistently realized in this scenario for a broad parametric range and in excellent agreement with the Planck legacy data. Distributed mass models can be viewed as realizations of the landscape property of string theory, with the mass distributions coming from the underlying spectra of the theory, which themselves would be affected by the vacuum of the theory. We discuss the recently proposed swampland criteria for inflation models on the landscape and analyze the conditions under which they can be met within the distributed mass warm inflation scenario. We demonstrate mass distribution models with a range of consistency with the swampland criteria including cases in excellent consistency.

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I. INTRODUCTION

The most recent cosmological observations once again confirm an expanding universe that is spatially flat, homogeneous and isotropic on large scales, and where the large scale structure originated from primordial fluctuations with a nearly scale-invariant, adiabatic, and gaussian spectrum [1]. Inflation [2] remains the dominant paradigm that can consistently explain the observational data. In the standard inflation picture, cold inflation (CI), depicted by a homogeneous scalar “inflaton” field, the short period of quasi-de Sitter accelerated expansion phase quickly dilutes away all traces of any pre-inflationary matter or radiation density, so that the state of the universe is the vacuum state. However, this generates a supercooled universe and leaves indeterminate a reasonable description of the transition from inflation to the “hot Big Bang” scenario, required by Big Bang Nucleosynthesis, and the physics of recombination leading to the Cosmic Microwave Background (CMB) that we observe today. This necessarily requires the conversion of inflaton energy density into ordinary matter and radiation and thus to its interactions with other fields.

In the conventional inflation picture the inflaton decay can only play a significant role at the end of the slow-roll regime, since particle production is not pictured to occur within the inflationary expansion phase. This leads to cold inflation ending through the standard “(p)reheating” paradigm [3]. The reasoning behind this phase lies in the fact that the perturbative decay width of a particle is generically smaller than its mass, which in turn lies below the Hubble expansion rate for a slowly rolling scalar field. Consequently such interplay between the inflaton and other constituents may perform a negligible role during the slow-roll phase of inflationary models. Nonetheless, it is relevant to note that the perturbative decay width only describes the decay of a field close to the minimum of its potential [4], which is evidently not the case during slow-roll dynamics, and that finite temperature effects can further significantly enhance the rate at which the inflaton dissipates its energy into other degrees of freedom.

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Thereby the inflaton field could be coupled to other components and might dissipate its vacuum energy and warm up the universe. This alternative scenario is known as the *warm inflation* (WI) paradigm [5, 6], where dissipative effects and associated particle production can, in fact, sustain a thermal bath concurrently with the accelerated expansion of the universe during inflation.

One of the earliest warm inflation models [7] suggested the idea that parameters in an inflation model could be randomly distributed. The distributed-mass-model (DM model) [8–11] was subsequently proposed and built on this idea in the context of string theory. It observed [9] that models from string theory have states at many energy levels and the distribution of these levels is ultimately dictated by the string vacuum. In particular, depending on the details of the compactification, the state of Kaluza-Klein modes, and patterns of symmetry breaking, it will induce splittings of string energy levels. Thus, based on the specific properties of a given string realization, the string states will be distributed differently. The idea of such distribution of states can find a more dynamical explanation in the context of the landscape picture, whereby different ground states of string theory will in turn imply changes to the spectrum of string states. In this respect, the DM model was one of the first low-energy realizations following the landscape idea well before the idea was formally stated.

The string landscape idea [12, 13] claims an unimaginable huge number of possible vacua reaching by some estimates to order 10^{500} . To have any hope in dealing with such a large number of possibilities, one would need to work both from the direction of string theory and phenomenology to identify viable vacua. The string landscape emerges from the complex structure of string theory. One property of this complex structure is the huge number of string states, going into the hundreds of thousands. Inflation models motivated by the landscape are generally built with only a small number of fields and generally do not attempt to utilize the vast spectra of string states. DM models are one of the few that attempt to utilize as part of the dynamical solution this inherent feature that there are many string states available. In the landscape way of thinking therefore, DM models, although phenomenologically motivated, presumably are studying theories emerging from a vast range of vacua that otherwise are missed in inflation models with just a small number of fields. As such, in constructing DM models it is natural to look for distributions that are of observational interest. In this paper we will first look at the generic DM distribution, where the string states in the range relevant to inflation are equally spaced apart, which was the original example studied in the early work [8–11]. This case has many relevant features but we find that it is not consistent with present observational constraints. We will then examine other types of DM distribution that have appealing consequences in comparison with observation.

Inflation is assumed to be described by low-energy Effective Field Theory (EFT), although in many models inflation can happen when the inflaton field is super-Planckian, particularly by considering monomial chaotic potentials in cold inflation. Furthermore, an EFT can be ultraviolet (UV) complete if it can be successfully incorporated in a quantum theory of gravity, such as String Theory. This paradigm provides a vast landscape of consistent embeddings of the EFT of gravity into a quantum theory, but this does not imply that any EFT coupled to gravity is consequently included in the landscapes. Those EFT's that are in fact inconsistent with a quantum theory of gravity lie in the surrounding swamplands [14, 15]. Hence, a benchmark is needed to ensure a de Sitter vacuum EFT can live in the desired string landscapes. Recently two swampland criteria relevant for inflationary theories have been proposed [16, 17], $|\Delta\phi|/M_P < \Delta$ and $M_P|V_\phi|/V > c$, provided that $V > 0$, where $\{\Delta, c\} \sim \mathcal{O}(1)$. Nonetheless, these criteria have been noted to pose inherent threats to the basic mechanism of slow-roll in cold inflation [17]. However, as part of the subsequent analysis we evaluate the aforementioned criteria in the warm inflation scenario. We show models of warm inflation that can very consistent with the swampland criteria. Inherently the dissipative feature in warm inflation makes it amenable for consistency with swampland criteria, as already noted in the literature [18–22].

In this paper we will examine in detail DM models and explore their consistency with observational data and theoretical viability. This work is organized as follows. In Sect. II we introduce all WI dynamics and primordial perturbation spectrum. In Sect. III we present in detail the supersymmetric distributed mass model, which is described by the interactions between the inflaton field and other light constituents: fermion ($\{\psi_i, \psi_\sigma\}$) and scalar ($\{\chi_i, \sigma\}$) fields. In Sect. IV we calculate all relevant parameters of the dissipation dynamics for the bosonic and fermionic sector: dissipative coefficients and their corresponding thermal average of the decay width. In Sect. V we develop a form of the mass distribution function. In Sect. VI we analyze in detail various dissipative coefficients for several inflation driven monomial potentials; we apply standard slow-roll methods and identify observationally consistent regions in parameter space. Here we also test the recently proposed swampland criteria. The main conclusions of this work are summarized in Sect. VII. Three appendices are also included, where we provide more detailed discussions of some of the results used in our computations.

II. WARM INFLATION DYNAMICS AND PRIMORDIAL PERTURBATION SPECTRUM

Non-equilibrium effects in the dynamics of a scalar field are generically produced due to interactions with an ambient thermal bath. The leading non-equilibrium effect, for a field evolving slowly compared to the characteristic time scale of the thermal bath, is a dissipative friction term $\Upsilon\dot{\phi}$ in its equation of motion [23], where $\Upsilon = \Upsilon(\phi, T)$ can be computed from first principles given the form of the interactions between the scalar field and the thermalized degrees of freedom. For a homogeneous field, this implies the continuity equation $\dot{\rho}_\phi + 3H(\rho_\phi + p_\phi) = -\Upsilon\dot{\phi}^2$, such that overall energy-momentum conservation implies the existence of an identical term with opposite sign in the continuity equation for the thermal fluid. This explicitly proves that dissipative effects in the inflaton's equation of motion lead to particle production in the thermal bath, so it prevents the exponential dilution of the latter in a quasi-de Sitter background. Therefore the temperature does not drop abruptly and the universe is able to smoothly cross to the radiation epoch. Taking into account dissipative effects, the evolution equation for the background inflaton field is given by:

$$\ddot{\phi} + (3H + \Upsilon)\dot{\phi} + V_{eff,\phi}(\phi, T) = 0, \quad (1)$$

where $V_{eff,\phi} = dV_{eff}/d\phi = V(\phi)_{,\phi} + V_{T,\phi}$ includes the effective thermal potential $V_{eff} = V(\phi) + V_T$, where $V(\phi)$ is the zero temperature potential, and V_T is the finite temperature effective potential, including radiative corrections to the potential due to any light component (fermion or scalar fields); Υ is the dissipative coefficient and H is the Hubble parameter, given by the Friedmann equation for a flat FRW universe:

$$H^2 = \frac{\rho_T}{3M_P^2}, \quad (2)$$

where $\rho_T = \dot{\phi}^2/2 + V_{eff} + Ts_R$ is the total energy density of the system, s_R is the entropy density, and M_P is the reduced Planck mass. One parameter that quantifies dissipation during WI is defined as the dissipative ratio $Q = \Upsilon/(3H)$. Depending on the ratio Q we can have different regimes: when $Q < 1$, this is called weak dissipative warm inflation; and when $Q \gtrsim 1$ we are in strong dissipative warm inflation. Furthermore, the WI paradigm assumes the presence of a thermal bath, at temperature T ; hence the evolution of such dissipative mechanism can be obtained from the evolution equation for the entropy density, given by:

$$T(\dot{s}_R + 3Hs_R) = \Upsilon\dot{\phi}^2, \quad (3)$$

which in the slow-roll regime reduces to $3HTs_R \simeq \Upsilon\dot{\phi}^2$. Without including the T -dependent corrections in the inflaton potential, we would have the standard relation $Ts_R = 4\rho_R/3 = 4C_RT^4/3$, where $C_R = \pi^2 g_{eff}(T)/30$, $g_{eff}(T)$ is the effective T -dependent number of relativistic degrees of freedom, with ρ_R denoting the standard energy density of radiation. All relativistic light fields contribute to the effective degrees of freedom, having:

$$g_{eff}(T) = \frac{90}{4\pi^2} \frac{s_R}{T^3}. \quad (4)$$

Since the radiative corrections modify the potential and its derivatives, they alter the standard slow-roll parameters, being from $\epsilon_\phi = M_P^2(V_{,\phi}/V)^2/2$ and $\eta_\phi = M_P^2 V_{,\phi\phi}/V$ to:

$$\epsilon_{eff} = \frac{M_P^2}{2} \left(\frac{V_{eff,\phi}}{V_{eff}} \right)^2, \quad \eta_{eff} = M_P^2 \frac{V_{eff,\phi\phi}}{V_{eff}}, \quad (5)$$

where $V_{eff,\phi\phi} = V(\phi)_{,\phi\phi} + V_{T,\phi\phi}$. Recall that inflation happens when $\dot{\phi}^2/2 \ll V_{eff}$ so does $(Ts_R Q^{-1}/2) \ll V_{eff}$, but even if small compared to the inflaton effective potential, it can be larger than the expansion rate with $(Ts_R)^{1/4} \gtrsim H$; by assuming thermalization, this translates roughly into $T \gtrsim H$, so one can obtain a warm inflationary universe consistently with a slow-roll evolution, as long as the radiative corrections are restrained such they do not spoil inflation. Furthermore, one can show that the radiation energy density portrayed by an entropic description can never exceed the inflationary potential in a slow-roll regime, guaranteeing a period of accelerated expansion:

$$\frac{Ts_R}{V_{eff}} \simeq \frac{2}{3} \frac{\epsilon_{eff}}{1+Q} \frac{Q}{1+Q}, \quad (6)$$

such that consistency of the slow-roll evolution requires $\epsilon_{eff} < 1 + Q$. This in turn also implies that, at the end of the slow-roll regime, when $\epsilon_{eff} \sim 1 + Q$, one may attain $T_{SR} \sim V_{eff}$ if a strong dissipative regime $Q \gtrsim 1$ can be achieved. In such cases radiation will smoothly become the dominant component at the end of inflation, providing the necessary “graceful exit” into the “hot Big Bang” cosmic evolution [24]. Although there may be additional particle production at the end of inflation, no reheating is actually necessary in WI when strong dissipation is reached, otherwise such mechanism is needed. Additionally to the smooth exit from inflation, WI exhibits several attractive features that have been explored in recent years. For instance, the dissipative friction damps the inflaton’s evolution, making slow-roll easier or, equivalently, alleviating the conditions on the flatness of the inflaton potential, expressed now by the slow-roll conditions $\epsilon_{eff}, |\eta_{eff}| \ll 1 + Q$. This may potentially provide a solution to the so-called “eta-problem” typically found in string/supergravity inflationary models where generically $\eta_\phi \sim \mathcal{O}(1)$ [10, 25]. (For other recent reviews of warm inflation, please see [26, 27].)

Small fluctuations of the inflaton about its homogenous component provide the initial seeds of density perturbation. These density perturbations produced during inflation evolve into the classical inhomogeneities observed in the CMB. For WI scenarios the fluctuations of the inflaton are thermally induced. As such, these initial seeds of density perturbations are already classical upon definition. The general expression for the amplitude of the primordial spectrum is given by [4, 28–31]:

$$\Delta_{\mathcal{R}}^2 = \left(\frac{H_*}{\dot{\phi}_*}\right)^2 \left(\frac{H_*}{2\pi}\right)^2 \left(1 + 2n_* + \frac{2\sqrt{3}\pi Q_*}{\sqrt{3 + 4\pi Q_*}} \frac{T}{H}\right) G(Q_*), \quad (7)$$

where all quantities are evaluated when the relevant CMB modes become superhorizon 50-60 e-folds before inflation ends. In the expression above, n_* denotes the inflaton phase space distribution at horizon crossing. By the strength of the interactions between the inflaton field and other particles in the thermal bath (including e.g. scattering processes), this might interpolate between the Bunch-Davies vacuum, $n_* = 0$, and the Bose-Einstein distribution at the ambient temperature T , $n_* \simeq (e^{H_*/T_*} - 1)^{-1}$. We will focus on the latter limiting case in this paper, which we denote as “thermal” inflaton fluctuations. Also the function $G(Q_*)$ accounts for the growth of inflaton fluctuations due to the coupling to radiation fluctuations through the temperature dependence of the dissipation coefficient and must be determined numerically. In addition, this function also exhibits a mild dependence on the form of the scalar potential. In general, with thermalised inflation fluctuations $1 + 2n_* = \coth(H_*/(2T_*))$, we have:

$$\Delta_{\mathcal{R}}^2 \simeq \left(\frac{3H_*^3(1 + Q_*)}{2\pi V_{eff,\phi_*}}\right)^2 \left(\frac{2\sqrt{3}\pi Q_*}{\sqrt{3 + 4\pi Q_*}} \frac{T_*}{H_*} + \coth\left(\frac{H_*}{2T_*}\right)\right) G(Q_*), \quad (8)$$

then from the amplitude of the curvature power spectrum, we may determine the scalar spectral index $n_s - 1 \simeq d \ln \Delta_{\mathcal{R}}^2 / dN_e$. Since, for $T \ll M_P$, gravitational waves are not significantly affected by thermal effects, the primordial tensor spectrum is given by the standard inflationary form $\Delta_t^2 = 2H_*^2/(\pi^2 M_P^2)$. The tensor-to-scalar ratio $r = \Delta_t^2/\Delta_{\mathcal{R}}^2$ is nevertheless affected, and in fact typically reduced, by the modifications to the scalar curvature perturbations introduced due to dissipation, which are basically a function of T_*/H_* and Q_* . We illustrate this fact by using the slow-roll dynamics, where the ratio r can be written as:

$$r \simeq \frac{16\epsilon_{eff}}{(1 + Q_*)^2 F(T_*/H_*, Q_*)}, \quad F(T_*/H_*, Q_*) = \left(\frac{2\sqrt{3}\pi Q_*}{\sqrt{3 + 4\pi Q_*}} \frac{T_*}{H_*} + \coth\left(\frac{H_*}{2T_*}\right)\right) G(Q_*), \quad (9)$$

which is suppressed w.r.t. the CI prediction by a factor $(1 + Q_*)^2 K(T_*/H_*, Q_*) > 1$. Indeed in [25] was shown explicitly, even before the BICEP and Planck results, that the presence of radiation and dissipation suppresses the tensor-to-scalar ratio. Authors in [25] computed the tensor-to-scalar ratio of the monomial ϕ^2 and ϕ^4 models and was one of the few analyses at the time that predicted for these inflation driven potentials a low tensor-to-scalar ratio, which now we see is consistent with data. Subsequent work further developed the analysis [32–36], showing that the tensor-to-scalar ratio may attain values even below 10^{-3} for a ϕ^4 potential, and thus potentially distinguishable from e.g. scenarios with a non-minimal coupling to gravity such as Higgs inflation [37] or the Starobinsky model [2].

III. SUPERSYMMETRIC DISTRIBUTED MASS MODEL

Let us consider the general form of an effective $N = 1$ global SUSY theory version of the distributed mass (DM) model with chiral superfields Φ , X_i and Y_i , described by the superpotential [25, 38–40]:

$$W = \sum_i \left[\frac{g}{2} (\Phi - M_i) X_i^2 + \frac{h}{2} X_i Y_i^2 \right], \quad (10)$$

where g and h are coupling constants and the sum is taken over an arbitrary distribution of supermultiplets X_i and Y_i . The chiral superfields Φ , X_i and Y_i have (scalar, fermion) components (ϕ, ψ_ϕ) , (χ_i, ψ_{χ_i}) and $(\sigma_i, \psi_{\sigma_i})$, respectively. Note that these are complex scalars and Weyl fermions, each with two degrees of freedom. We may use the Majorana representation for the spinors, i.e. use a 4-component Majorana spinor built from the same Weyl fermion. In this case, note that a Majorana fermion is its own anti-particle. We have considered different Y_i fields coupled to each X_i field in the tower to avoid mass mixing at the level of the thermal masses. The scalar interaction terms in the theory are obtained from the superpotential using [38]:

$$-\mathcal{L}_S = |\partial_\Phi W|^2 + \sum_i |\partial_{X_i} W|^2 + \sum_i |\partial_{Y_i} W|^2. \quad (11)$$

On the other hand the fermion Lagrangian can be computed from the general formula [38]:

$$-\mathcal{L}_F = \frac{1}{2} \sum_{n,m} \frac{\partial^2 W}{\partial \xi_n \partial \xi_m} \bar{\psi}_n P_L \psi_m + \frac{1}{2} \sum_{n,m} \frac{\partial^2 W^\dagger}{\partial \xi_n^\dagger \partial \xi_m^\dagger} \bar{\psi}_n P_R \psi_m, \quad (12)$$

where ξ_n is a superfield: Φ, X_i, Y_i , and $P_L = 1 - P_R = (1 + \gamma_5)/2$ are the chiral projection operators acting on Majorana 4-spinors. Note that $\partial W / \partial X_i \partial Y_i = \partial W / \partial Y_i \partial X_i$ and similarly $\partial W^\dagger / \partial X_i^\dagger \partial Y_i^\dagger = \partial W^\dagger / \partial Y_i^\dagger \partial X_i^\dagger$. Then we select only the boson field components of the chiral superfields Φ , X_i and Y_i , which are ϕ , χ_i and σ_i , respectively. Nonetheless, see that the inflaton field ϕ can be decomposed into its real and imaginary parts via^a $\phi = (\phi_R + i\phi_I)/\sqrt{2}$. Hence, this prescription introduces another decay channel; for instance, the modulus square becomes $|\phi - M_i|^2 = (\phi_R/\sqrt{2} - M_i)^2 + \phi_I^2/2$. However, $\phi_R/\sqrt{2}$ is the nonzero vacuum expectation value of ϕ , which will be the only term that will contribute to dissipation in the scalar field's effective equation of motion, so that the imaginary part of the inflaton, ϕ_I , is not relevant for our subsequent analysis. Moreover, in order to be coherent with further calculations, ϕ is going to be considered only as the classical expectation value, without the label R and the factor $1/\sqrt{2}$. Therefore, the relevant Lagrangians that may contribute to dissipation in the scalar field's effective equation of motion are:

$$\begin{aligned} -\mathcal{L}_S &= g^2 \sum_i (\phi - M_i)^2 |\chi_i|^2 + \frac{gh}{2} \sum_i (\phi - M_i) \left[\chi_i (\sigma_i^\dagger)^2 + \chi_i^\dagger \sigma_i^2 \right] \\ &\quad + h^2 \sum_i |\chi_i|^2 |\sigma_i|^2 + \frac{g^2}{4} \sum_i |\chi_i|^4 + \frac{h^2}{4} \sum_i |\sigma_i|^4, \end{aligned} \quad (13)$$

$$\begin{aligned} -\mathcal{L}_F &= \frac{g}{2} \sum_i (\phi - M_i) \bar{\psi}_{\chi_i} P_L \psi_{\chi_i} + \frac{g}{2} \sum_i (\phi - M_i) \bar{\psi}_{\chi_i} P_R \psi_{\chi_i} \\ &\quad + \frac{h}{2} \sum_i \chi_i \bar{\psi}_{\sigma_i} P_L \psi_{\sigma_i} + \frac{h}{2} \sum_i \chi_i^\dagger \bar{\psi}_{\sigma_i} P_R \psi_{\sigma_i} + h \sum_i \sigma_i \bar{\psi}_{\sigma_i} P_L \psi_{\chi_i} + h \sum_i \sigma_i^\dagger \bar{\psi}_{\sigma_i} P_R \psi_{\chi_i}. \end{aligned} \quad (14)$$

Note that the bare masses are $m_{\psi_{\chi_i}} = g(\phi - M_i) = m_{\chi_i}$ at zero temperature for unbroken SUSY^b. This is in agreement with [38] upon rescaling the couplings g and h in the superpotential by $1/2$ factors.

^a Similarly we define the complex scalars in terms of the real and imaginary part for χ_i and σ_i .

^b Note that in warm inflation SUSY is broken both by the finite temperature and the inflaton energy density. The latter should arise from an additional Φ -dependent term in the superpotential that we have not included above and that will lead to a small splitting of the mass for the real and imaginary components of the χ_i scalar fields.

At finite temperature, both the χ_i and ψ_{χ_i} receive thermal mass corrections. The contributions of the σ_i and ψ_{σ_i} fields to the latter have been computed in [38] and been shown to be identical for both χ_i and ψ_{χ_i} , corresponding to $h^2 T^2/8$ taking into account the coupling normalization differences. The χ_i scalars also receive thermal corrections from their self-interactions $g^2 |\chi_i|^4/4$. Noting that these interactions give a contribution to their tree-level mass $\partial^2 V/\partial \chi_i \partial \chi_i^\dagger = g^2 |\chi_i|^2$ and taking into account the contribution of the χ_i fields to the thermal effective potential:

$$\Delta V_T \subset 2 \times \frac{m_{\chi_i}^2}{24} T^2 = \frac{g^2}{12} T^2 |\chi_i|^2 + \dots, \quad (15)$$

where we have taken into account the two degrees of freedom for complex scalars, this yields a thermal mass correction $g^2 T^2/12$ to the χ_i fields^c. In summary, we obtain:

$$\Delta m_{\chi_i}^2 = \frac{g^2}{12} T^2 + \frac{h^2}{8} T^2, \quad \Delta m_{\psi_{\chi_i}}^2 = \frac{h^2}{8} T^2. \quad (16)$$

We also need to compute the finite temperature decay widths of the χ_i and ψ_{χ_i} fields. However, we only recall such calculations for Dirac fermions [39]; albeit the difference between Majorana and Dirac fermions is only in the overall factors of the decay width. Hence, we can first compute them at zero temperature to set the correct normalisation factors in order to identify these global constants. In general we have for the decay of a particle of mass m at rest into a pair of massless particles:

$$\Gamma = \frac{S}{16\pi m} |\mathcal{M}|^2, \quad (17)$$

where $S = 1/2$ if the particles are identical and $S = 1$ if they are distinct. The χ_i fields may decay via $\chi_i \rightarrow \sigma_i \sigma_i$ and $\chi_i \rightarrow \psi_{\sigma_i} \psi_{\sigma_i}$. In the first case, dropping the indices for simplicity, we may decompose the fields into their real and imaginary components via $\chi = (\chi_R + i\chi_I)/\sqrt{2}$ and analogously for σ . This yields the scalar interactions:

$$-\mathcal{L}_{\chi\sigma^2} = \frac{hg}{2\sqrt{2}} (\phi - M_i) [\chi_R \sigma_R^2 - \chi_R \sigma_I^2 + 2\chi_I \sigma_R \sigma_I]. \quad (18)$$

Hence, the χ_R scalar may decay into σ_R or σ_I pairs, while the χ_I scalar has only one decay channel $\chi_I \rightarrow \sigma_R \sigma_I$. For each of these decay channels, the vertex factor is $-ihg(\phi - M_i)/\sqrt{2}$. Taking into account the $S = 1/2$ factors in the $\chi_R \rightarrow \sigma_R \sigma_R$ and $\chi_R \rightarrow \sigma_I \sigma_I$ channels, we then find that the decay widths are equal for both χ_R and χ_I , being given by:

$$\Gamma_{\chi_i}^S = \frac{h^2 g^2 (\phi - M_i)^2}{32\pi m_{\chi_i}}. \quad (19)$$

The fermionic decay channel comes from the interaction term $\frac{1}{2} h \chi_i \bar{\psi}_{\sigma_i} \psi_{\sigma_i}$, where the vertex factor is simply $-ih$. Since the particles are identical in the final state and computing the matrix element with the usual Feynman rules, this yields:

$$\Gamma_{\chi_i}^F = \frac{h^2}{32\pi} m_{\chi_i}. \quad (20)$$

Noting that, at zero temperature, $m_{\chi_i} = g(\phi - M_i)$, we see that $\Gamma_{\chi_i}^S = \Gamma_{\chi_i}^F$ in this limit.

The fermions ψ_{χ_i} can decay as $\psi_{\chi_i} \rightarrow \sigma_i \psi_{\sigma_i}$ via the corresponding Yukawa term above, and this yields simply:

$$\Gamma_{\psi_{\chi_i}} = \frac{h^2}{16\pi} m_{\psi_{\chi_i}}. \quad (21)$$

^c The mass of the scalar σ_i fields is also corrected by a similar factor (involving only the h coupling) and, although the resulting masses are below the temperature, they will generically lie above H in the parametric range relevant to our discussion. This implies that no scalar field other than the inflaton can sustain a slowly-rolling background value that drives inflation.

Again, note that at zero temperature for unbroken SUSY we have $m_{\chi_i} = m_{\psi_{\chi_i}} = g(\phi - M_i)$, which yields identical total decay widths for the scalars and fermionic superpartners, as it should.

Once we have compute the decay widths at zero temperature, we identify each vertex factor squared from the decay of a scalar boson to fermions and the decay of a fermion to a scalar boson and a fermion. In the pole approximation (see Appendix A) we take the limit $T \rightarrow 0$ of the corresponding decay widths [41] and then by comparing them with eqs. (20) and (21), we identify: $g_{\chi_i \bar{\psi}_{\sigma_i} \psi_{\sigma_i}}^2 = h^2/4$ and $g_{\psi_{\chi_i} \sigma_i \psi_{\sigma_i}}^2 = h^2$. From here we will use these overall factors for the rest of the computations.

IV. DISSIPATIVE INTERACTIONS

The particle physics model considered in this work is inspired by string theory exhibiting $\mathcal{N} = 1$ global supersymmetry, with the inflaton field coupled to massive modes of the string, as discussed in [42]. Several interactions are identified by the shifted couplings $g^2(\phi - M_i)^2 \chi_i^2$, and $g(\phi - M_i) \bar{\psi}_{\chi_i} \psi_{\chi_i}$ for bosons χ_i and fermions ψ_{χ_i} respectively, with $\{M_i\}$ ranging over mass scales. This feature yields the name of distributed-mass-model (DM-model). In the subsequent segments we establish the relevant interplay the inflaton field has with the aforementioned fields. We will restrict to interactions such that the leading contribution will come dominantly from one-loop processes, when the decaying field is light. The key property of the DM model is that for a given temperature T , only the fields with masses $g^2(\phi - M_i)^2 \lesssim T^2$ will contribute to the dissipation. Henceforth, we will refer such configuration of states as thermally excited sites. We will consider separately the dissipative processes associated with the excitation of the scalar χ_i fields, which may decay via $\chi_i \rightarrow \sigma_i \sigma_i$ or $\chi_i \rightarrow \bar{\psi}_{\sigma_i} \psi_{\sigma_i}$, and those associated with the excitation of the fermionic ψ_{χ_i} fields, which decay via $\psi_{\chi_i} \rightarrow \sigma_i \psi_{\sigma_i}$, with technical details of the computation given in Appendix A.

A. Bosonic sector

The dissipative coefficient arising from the pattern of interactions among the scalar component and the light states is given by [8, 10, 11](see also Appendix A):

$$\Upsilon^S(\phi, T) = \sum_{i=1}^{t.e.} \frac{32g^4}{\pi h^2 \left[16 \frac{m_{\chi_i}^2}{\tilde{m}_{\chi_i}^2} + \frac{\tilde{m}_{\chi_i}^2}{T^2} \right]} \ln \left(\frac{2T}{\tilde{m}_{\chi_i}} \right) \frac{(\phi - M_i)^2}{\tilde{m}_{\chi_i}}, \quad (22)$$

where ‘t.e.’ means sum over all thermally excited sites. In the computation of the dissipative coefficient, the total decay rate for all processes is already taken into account, which is given by [8, 10, 11](see also Appendix A):

$$\Gamma_{\chi_i} = \Gamma(\chi_i \rightarrow \sigma_i \sigma_i) + \Gamma(\chi_i \rightarrow \bar{\psi}_{\sigma_i} \psi_{\sigma_i}) = \frac{h^2}{128\pi} \frac{T \tilde{m}_{\chi_i}}{\omega_{\chi_i}(p)} \left[16 \frac{m_{\chi_i}^2}{\tilde{m}_{\chi_i}^2} + \frac{\tilde{m}_{\chi_i}^2}{T^2} \right], \quad (23)$$

where $\omega_{\chi_i}(p) = \sqrt{\tilde{m}_{\chi_i}^2 + |\mathbf{p}|^2}$ and \tilde{m}_{χ_i} is the full χ_i mass, including the tree-level contribution m_{χ_i} and the thermal corrections computed above. We need to ensure that all light bosons remain in a nearly-thermal state during inflation, so that the above result for the dissipation coefficient is a consistent approximation. For simplicity, we may take the thermal average of the above decay width, such:

$$\bar{\Gamma}_{\chi_i} = \frac{1}{n_B} \int \frac{d^3 p}{(2\pi)^3} \Gamma_{\chi_i} f_B \quad (24)$$

where $f_B(\mathbf{p}/T)$ is the Bose-Einstein distribution and $n_B(\tilde{m}_{\chi_i}/T)$ the associated number density. This yields in the limit $m_{\chi_i} \lesssim T$ via a numerical approximation (see Appendix B):

$$\bar{\Gamma}_{\chi_i} \simeq \frac{h^2 f^{1/2}}{128\pi} \left[\frac{16 m_{\chi_i}^2}{f T^2} + f \right] \frac{0.68}{(1 + 0.77 f^{1/2})} T, \quad (25)$$

where $f = f(g, h) = g^2/12 + h^2/8$. Note that the above expression depends on the factor $m_{\chi_i}/T = g(\phi - M_i)/T$, so that different fields will decay at different rates. However, note that states for which $m_{\chi_i}/T = 0$, corresponding to $\phi = M_i$, do not contribute to dissipation according to Eq. (22). Conversely, the heaviest states that can be thermally excited have $m_{\chi_i} \sim T$, and these are the ones that contribute the most to dissipation at any given time. Therefore, in analyzing the consistency of the model, namely whether the χ_i decay faster than expansion in order to remain close to thermal equilibrium, we will consider the states for which $m_{\chi_i}/T \simeq 1$.

B. Fermionic sector

The calculation of the dissipation coefficient can be done following essentially the same steps as in the ‘‘Warm Little Inflaton’’ model [34–36], yielding:

$$\Upsilon^F(T) = \sum_{i=1}^{t.e.} C_T^F T, \quad C_T^F \simeq \frac{3g^2}{h^2(1 - 0.34 \log(h))}. \quad (26)$$

This calculation involves computing the finite temperature decay width of all interaction processes, which is given by

$$\Gamma_{\psi_{\chi_i}} = \Gamma(\psi_{\chi_i} \rightarrow \sigma_i \psi_{\sigma_i}) = \frac{h^2}{16\pi} \frac{T^2 m_{\psi_i}^2}{\omega_{\psi_i}^2(p) |\mathbf{p}|} \left[F\left(\frac{k_+}{T}, \frac{\omega_{\psi_i}(p)}{T}\right) - F\left(\frac{k_-}{T}, \frac{\omega_{\psi_i}(p)}{T}\right) \right], \quad (27)$$

where, neglecting the masses of the decay products σ and ψ_σ , we have $\omega_{\psi_i}(p) = \sqrt{\tilde{m}_{\psi_i}^2 + |\mathbf{p}|^2}$, $k_\pm = (\omega_{\psi_i}(p) \pm |\mathbf{p}|)/2$ and

$$F(x, y) = xy - \frac{x^2}{2} + (y - x) \ln \left(\frac{1 - e^{-x}}{1 + e^{x-y}} \right) + \text{Li}_2(e^{-x}) + \text{Li}_2(-e^{x-y}), \quad (28)$$

where $\text{Li}_2(z)$ is the dilogarithm function. Note that the mass of the light fields is corrected by a factor $h^2 T^2/8$ due to their interactions with the thermal bath fields ψ_{σ_i} and σ_i , which we will take to be dominant over the inflaton contribution, i.e. we take $\tilde{m}_{\psi_i} \simeq h^2 T^2/8$ in the above computation of the decay width. We need to ensure that all light fermions remain in a nearly-thermal state during inflation, so that the above result for the dissipation coefficient is a consistent approximation. For simplicity, we may take the thermal average of the above decay width in the high-temperature and ultra-relativistic region:

$$\bar{\Gamma}_{\psi_{\chi_i}} = \frac{1}{n_F} \int \frac{d^3 p}{(2\pi)^3} \Gamma_{\psi_{\chi_i}} f_F \quad (29)$$

where $f_F(\mathbf{p}/T)$ is the Fermi-Dirac distribution and $n_F(\tilde{m}_{\chi_i}/T)$ the associated number density. This yields in the limit $m_{\psi_{\chi_i}} \lesssim T$:

$$\bar{\Gamma}_{\psi_{\chi_i}} \simeq 10^{-3} \left[1 - 0.875 \ln \left(\frac{h^2}{8} \right) \right] h^4 T. \quad (30)$$

V. MASS DISTRIBUTION FUNCTION

In the DM model the dissipative mechanism is well described by the interaction between the inflaton ϕ and the finite number of fermion and boson fields that are light at any given time as the scalar field scans the tower of states. DM models are distinguished by how the mass sites are distributed. Such an idea has a natural realization with string theory, whereby the inflaton is suggestive of an excited string zero mode, which then interacts with massive string levels. Such a construction of DM models from string theory was shown in [9]. String levels can be highly degenerate and the distribution of mass states was suggested in [9] to emerge as a fine structure splitting of such

levels. Thus the pattern of splitting of a level will depend on the properties of the string state, ultimately governed by the underlying string vacuum. In this respect different string vacua should imply different distribution of mass states. A first principles determination of such distributions is difficult and minimally requires a detailed study of string theory case-by-case for each different possible vacua. However, we can adopt a phenomenological approach and look at various types of mass distributions and see what type of inflation they can lead to. This would be a minimum first step to test the viability of such models. In this section we develop some basic properties of the mass distribution function.

In general the mass of states labeled by i is as discussed in Sect. III $m_{\chi_i}^2 = m_{\psi_{\chi_i}}^2 = g^2(\phi - M_i)^2$ and so is governed by the parameter M_i . It is the distribution of these mass sites M_i that the underlying theory should determine, but here we treat them as phenomenological parameters and examine different types of distributions of M_i over mass sites labelled by i . To understand the meaning of such mass sites, suppose the inflaton field ϕ sat in the middle of some mass sites. Thus for all mass sites $\phi - T/g < M_i < \phi + T/g$, the corresponding fields χ_i and ψ_{χ_i} would be thermally excited. As some terminology, we will call the region $\Delta\phi$ surrounding the inflaton ϕ that has thermally excited fields, a thermal interval. Some examples of this are shown in Fig. 1. The idea of warm inflation in such models is that the inflaton rolls through a region with many such mass sites, thus thermally exciting for some time a given field and then once ϕ is far enough away, that field again is no longer thermally excited. This implies as the inflaton slow-rolls during inflation, a thermal interval surrounding it moves with it.

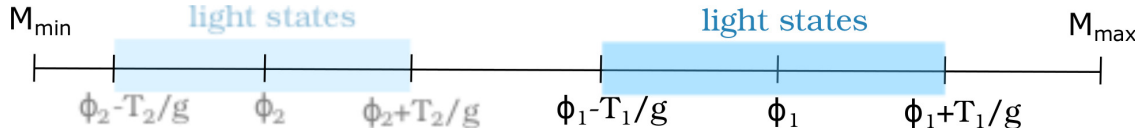


FIG. 1: Schematic representation of the time evolution in the model assuming that the temperature decreases.

With no underlying theory to dictate such mass distributions, we will simply consider one type of construction to make the idea more tangible. We will assume the mass sites can be written as

$$M_i = \phi + i\Delta M(\phi, T, m, g), \quad (31)$$

where ΔM is the mass gap in the tower, which may in general be a dynamical function of the inflaton's expectation value and the ambient temperature, as well as of some intrinsic mass scale m and couplings such as g . For instance, in a theory with compact extra-dimensions, as is the case of string/M-theory, the mass gap in the Kaluza-Klein tower is inversely proportional to the size of the extra-dimensions, generically set by the expectation value of one or more moduli fields. If such moduli interact with the inflaton and/or the thermal bath, their expectation value may become dynamical (although slowly evolving) during inflation, exhibiting a ϕ - and/or T -dependence^d. With this motivation in mind, we will take an effective field theory approach and consider functional forms of the mass gap that lead to simple forms for the dissipative coefficient, analyzing their dynamical and observational consequences.

In the expression above i is an integer ranging from mass sites closest to ϕ at $i = 1$ up to a maximum mass site that is thermally excited determined by $i_{max}\Delta M = T/g$. During the course of the observable range of inflation (around 50-60 e-folds), the inflaton will traverse some distance covering many thermal intervals $\Delta\phi$ and in general one needs mass distributions where $\Delta M \ll \Delta\phi$, so that the inflaton crosses many mass sites in each thermal interval. One expects at least tens of mass sites per thermal interval and hundreds to thousands of thermal intervals crossed during the observable period of inflation. In other words, in each e-fold of inflation one expects tens or few hundred mass sites to be crossed. Thus thousands of string states will be thermally excited for a brief period of time in the course of such a warm inflation period.

With this form of mass distribution function, one can then obtain expressions for the dissipative coefficient and

^d Note that the bare masses of the fields in the tower relevant for the dissipative process, $gM_i \sim g\phi$, are typically very large and that, at any given time, only a finite number of states become light due to the inflaton field reducing their mass. Therefore, this scenario does not require a low KK/string scale to be implemented in the context of string/M-theory.

the effective potential. For the dissipative coefficient the sum that is required is

$$\sum_{i=1}^{i_{max}} (\phi - M_i)^2 = \sum_{i=1}^{i_{max}} i^2 \Delta M^2 \approx i_{max}^3 \Delta M^2 \simeq T^3 / (g^3 \Delta M) . \quad (32)$$

For example if we wanted $\Upsilon \sim \phi$ it means considering $\Delta M \sim T/(g^3 \phi)$ and similarly for other forms of the dissipation function.

In the landscape context, the idea being there are a huge number of different vacua and some subset can produce a distributed mass model and some subset of those could have a distribution of mass sites that gives a dissipative coefficient that leads to observably consistent warm inflation. Since the possible vacua is so huge, its simply seen as a statistical possibility that somewhere in this landscape a mass distribution of a particular type can be found. One would need to see this in an anthropic way that there could be inflation of many different forms occurring over the landscape, but some will occur in a way favorable to create a Universe like ours. In this respect choosing a function ΔM boils down to finding the type of functions that can lead to a Universe that looks like ours. If that could successfully be realized, as we will examine in the next Section, then this model would have phenomenological viability. One could then explore as a much bigger step if its possible to find such particular types of DM models from string theory, although we will not be exploring that question here.

We can follow the above prescription in order to evaluate the effective finite temperature potential, where for the sake of simplicity we evaluate only the bare masses of the χ 's and ψ_χ 's fields. Hence both bosonic and fermionic sector, can be estimated using [51, 53, 54] (see eqs. (C7) and (C1)):

$$V_T^{\chi_i} \simeq \frac{T^5}{g\Delta M} \left\{ -\frac{\pi^2}{45} + \frac{1}{12} + \frac{1}{6\pi} - \frac{1}{32\pi^2} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_b \right] \right\} , \quad (33)$$

$$V_T^{\psi_{\chi_i}} \simeq \frac{T^5}{g\Delta M} \left\{ -\frac{7\pi^2}{360} + \frac{1}{24} + \frac{1}{32\pi^2} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right\} , \quad (34)$$

hence if we select $\Delta M \sim T/(g^3 \phi)$ the total effective finite temperature potential becomes:

$$V_T^{\chi_i} + V_T^{\psi_{\chi_i}} \simeq \frac{g^2 T^3}{24} \phi \left[-\pi^2 + 3 + \frac{4}{\pi} + \frac{3}{4\pi^2} (c_b - c_f) \right] . \quad (35)$$

Since the thermal corrections to the inflaton potential are, for this particular case, linear in the field, they contribute to the ϵ_{eff} parameter and need to be taken into account. However, in the scenarios with a quartic chaotic potential $V(\phi) = \lambda \phi^4$ that we analyze in more detail below, we typically find $T/\phi \sim 10^{-5}$ and $\lambda \sim 10^{-15} - 10^{-14}$, such that these thermal corrections give a negligible contribution to both the effective potential and its first derivative for any value of the coupling g in the perturbative regime. In addition, in the scenarios where the mass gap is field-independent, thermal corrections to the inflaton potential do not modify the slow-roll parameters.

VI. RESULTS

Let us now analyze the inflationary dynamics for a quartic scalar potential, $V(\phi) = \lambda \phi^4$, taking into account both scalar and fermionic contributions to the dissipation coefficient. In addition different types of mass distribution functions will be used leading to dissipative coefficients of various forms, namely $\Upsilon \propto T^2$, $\Upsilon \propto T$ and $\Upsilon \propto \phi$. Finite temperature corrections to the effective potential are computed for each case. These effects can be controlled thus preventing thermal effects from generating large contributions to the inflaton mass that could reintroduce the “ η -problem”.

A. $\Upsilon \propto T^2$

For a homogeneous distribution of states in the tower, with constant mass splitting ΔM , as proposed by [8, 10, 11], the resulting scalar coefficients grows with temperature as $\Upsilon \propto T^2$, with fermions contributing only about

20% to the total dissipation as first estimated in [8, 10, 11]. Due to the adiabatic condition $\bar{\Gamma}_\chi^S/H > 1$, one requires large values of T_*/H_* ; quantitatively we have that at least $T_*/H_* \gtrsim 150$, which in turn yields values of $Q_* \gtrsim 10$. Our results here are consistent with those reported in [8, 10, 11]. However in the strong dissipation regime it is now understood that the primordial perturbations have a growing mode [29, 30, 43, 44]. The consequences of this growing mode are to tilt the spectrum of perturbations towards the blue-tilted region, $n_s > 1$, and so inconsistent with the CMB data. Thus, the DM model with a homogeneous mass distribution resides outside the observational window provided by Planck data [1]. Although this particular mass distribution is not consistent with observation, the basic idea of the model shows appealing features and one can explore other types of mass distributions as we will now do.

$$\text{B. } \Upsilon = \tilde{C}_\phi \phi \text{ and } V(\phi) = \lambda \phi^4$$

We want to study this model within the strong dissipative regime, since we expect that this model fits perfectly in it. This facilitates all calculations, so we will use standard analytical tools, i.e. we may implement the standard slow-roll parameters $\epsilon_\phi = M_P^2(V_{,\phi}/V)^2/2$ and $\eta_\phi = M_P^2 V_{,\phi\phi}/V$. Also we introduce another slow-roll parameter to take into account the variation of Υ :

$$\beta = M_P^2 \frac{\Upsilon_{,\phi} V_{,\phi}}{\Upsilon V}, \quad (36)$$

and the slow-roll conditions are now given by $\epsilon_\phi < Q$, $|\eta_\phi| < Q$, $\beta_\phi < Q$. We use the slow-roll equations, at $Q \gg 1$, in order to find a direct relation between Q and ϕ , yielding:

$$Q = \frac{\tilde{C}_\phi}{\sqrt{3\lambda}} \frac{M_P}{\phi}, \quad (37)$$

and the slow-roll parameters are:

$$\epsilon_\phi = \frac{8M_P^2}{\phi^2}, \quad \eta_\phi = \frac{12M_P^2}{\phi^2}, \quad \beta_\phi = \frac{4M_P^2}{\phi^2} \quad (38)$$

so in the strong dissipative regime inflation ends at $\epsilon_\phi \simeq Q$, this yields the inflaton value:

$$\phi_{end} = \frac{8\sqrt{3\lambda}}{\tilde{C}_\phi} M_P. \quad (39)$$

The number of e-folds during inflation in the aforementioned regime is:

$$\begin{aligned} N_e &= -\frac{1}{M_P^2} \int_{\phi_*}^{\phi_{end}} d\phi \frac{QV}{V_\phi} = -\frac{\tilde{C}_\phi}{4\sqrt{3\lambda}M_P} \int_{\phi_*}^{\phi_{end}} d\phi \\ &= \frac{\tilde{C}_\phi}{4\sqrt{3\lambda}} \left[\left(\frac{\phi_*}{M_P} \right) - \left(\frac{\phi_{end}}{M_P} \right) \right], \end{aligned} \quad (40)$$

hence the inflaton value at horizon crossing is:

$$\phi_* = \frac{4\sqrt{3\lambda}(N_e + 2)}{\tilde{C}_\phi} M_P = 2\sqrt{\frac{N_e + 2}{Q_*}} M_P. \quad (41)$$

Once we have determined ϕ_* we can evaluate the observables at 50-60 e-folds before inflation ends. Before we proceed, one should note that in the strong dissipative regime many parameters and formulas simplify, in fact, one of them is an analytic approximation to the density perturbation amplitude. Recall that for $T > H$, the dominant contribution to the primordial perturbation spectrum are thermal fluctuations of the inflaton field, as opposed to the conventional quantum fluctuations in cold inflation models. Upon exiting the horizon these thermal fluctuations freeze out as

classical perturbations and during slow-roll at $Q \gg 1$ the amplitude of the curvature perturbation power spectrum is given by [28, 45]:

$$\Delta_{\mathcal{R}}^2 \simeq \frac{9}{4\pi^2} \frac{H_*^5 T_* Q_*^{5/2}}{V_{\phi_*}^2}, \quad (42)$$

where all quantities are evaluated at horizon-crossing. Hence the scalar spectral index $n_s - 1 \simeq d \ln \Delta_{\mathcal{R}}^2 / dN_e$ is modified as well, becoming [28]:

$$\begin{aligned} n_s &= 1 + \frac{3}{2} \left(\frac{\eta_V}{Q} - \frac{3}{2} \frac{\epsilon_V}{Q} - \frac{3}{2} \frac{\beta_V}{Q} \right) = 1 - \frac{9}{4} \left(\frac{2\sqrt{3}\lambda}{\tilde{C}_\phi} \frac{M_P}{\phi_*} \right) \\ &= 1 - \frac{9}{4} \left(\frac{1}{N_e + 2} \right). \end{aligned} \quad (43)$$

Remarkably, note that n_s only depend of the number of e-folds: $n_s(N_e = 50) = 0.9567$ and $n_s(N_e = 60) = 0.9637$; where at 60 e-folds this scalar spectral index agrees outstandingly with Planck data [1]. The tensor-to-scalar ratio $r = \Delta_t^2 / \Delta_{\mathcal{R}}^2$, as mentioned above, is typically reduced by the modifications to the scalar curvature perturbations introduced due to dissipation, which is basically a function of Q_* . We illustrate this fact by using the slow-roll dynamics, where the ratio r can be written as:

$$r = \frac{2H_*^2}{\pi^2 M_P^2 \Delta_{\mathcal{R}}^2} = \frac{2\lambda}{3\pi^2 \Delta_{\mathcal{R}}^2} \frac{\phi_*^4}{M_P^4} = \frac{32\lambda}{3\pi^2 \Delta_{\mathcal{R}}^2} \frac{(N_e + 2)^2}{Q^2} \quad (44)$$

The value of λ is fixed by using the normalization of amplitude of the primordial spectrum $\Delta_R^2 \simeq 2.2 \times 10^{-9}$ [1]. Using the slow-roll equations in Eq. (42), we have:

$$\lambda = (\Delta_{\mathcal{R}}^2)^{4/3} \left(\frac{(12\pi^4 C_R^{1/2})^{2/3}}{(N_e + 2)^3} \right), \quad (45)$$

where $C_R = \pi^2 g_{eff}/30$, and $g_{eff} = 1 + 15N_M/4$, N_M being the no of bosonic χ_i (fermionic ψ_i) light degrees of freedom at horizon crossing. The tensor-to-scalar ratio is then just a function of Q_* and N_e ,

$$r = \frac{64(6\Delta_{\mathcal{R}}^2)^{1/3}}{3Q_*^2(N_e + 2)}. \quad (46)$$

As one expected, the tensor-to-scalar ratio is highly suppressed by dissipation. In fact this ratio lies in the region $10^{-9} < r \lesssim 10^{-4}$ at $10 \leq Q_* \leq 1000$, for $N_M = O(10 - 100)$, having the bigger suppression at the largest Q_* .

We can also look at the ratios ϕ_*/M_P and m_{ϕ_*}/H_* , in order to illustrate that for this model the relevant scales can indeed happen in the sub-Planckian region, as shown in Fig. 2. This occurs thanks to the high dissipation dynamics, since in WI field potentials that are not flat enough to allow the standard slow-roll inflaton evolution, i.e. they do not have large enough initial inflaton values so not “sufficient” inflation takes place, can in fact lead to a longer period of inflation due to the extra friction induced by Υ . Hence, for small (even sub-Planckian) ϕ_* inflation can last 50-60 e-folds in the strong dissipative regime.

We can even generalise above results for a dissipative coefficient of the form: $\Upsilon = C_\phi \phi^p m^{1-p}$, where m is a mass scale and p is a free parameter. For this generic case, the scalar spectral index and the tensor-to-scalar ratio become very simple expressions that yield a remarkable agreement with Planck legacy [1] (see Fig. 3):

$$n_s = 1 - \frac{9}{4} \left(\frac{p}{pN_e + 2} \right), \quad r = \frac{32\lambda}{3\pi^2 \Delta_{\mathcal{R}}^2} \frac{(pN_e + 2)^2}{Q^2} \quad (47)$$

The good agreement between WI with a dissipative coefficient of the form: $\Upsilon = C_\phi \phi^p m^{1-p}$ described by a chaotic quartic potential has been reported before, for instance [30] analyzed the $\Upsilon \propto \phi^2$ case, finding observably favorable region for $Q \gg 1$. On the other hand, the tensor-to-scalar ratio is enhanced by larger p , so for very large p this ratio r can indeed become bigger than 1. The limits on ongoing and planned B-mode polarization experiments [46] suggest sensitivity down to $r \gtrsim 10^{-3}$. This leaves open the door to access desirable values of r for very large Q_* .

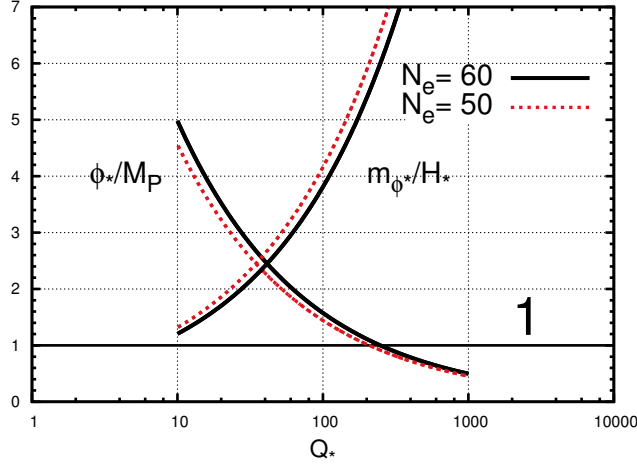


FIG. 2: Ratios m_{ϕ_*}/H_* and ϕ_*/M_P , as a function of the dissipative ratio Q_* , within a DM scenario with a dissipative coefficient of the form $\Upsilon \propto \phi$ described by a quartic potential for 50 (red dashed-line) and 60 (black solid-line) e-folds of inflation. Note that m_{ϕ_*}/H_* is always larger than 1 for $Q_* > 10$, while $\phi_*/M_P < 1$ (sub-Planckian) at $Q_* \sim 210$.

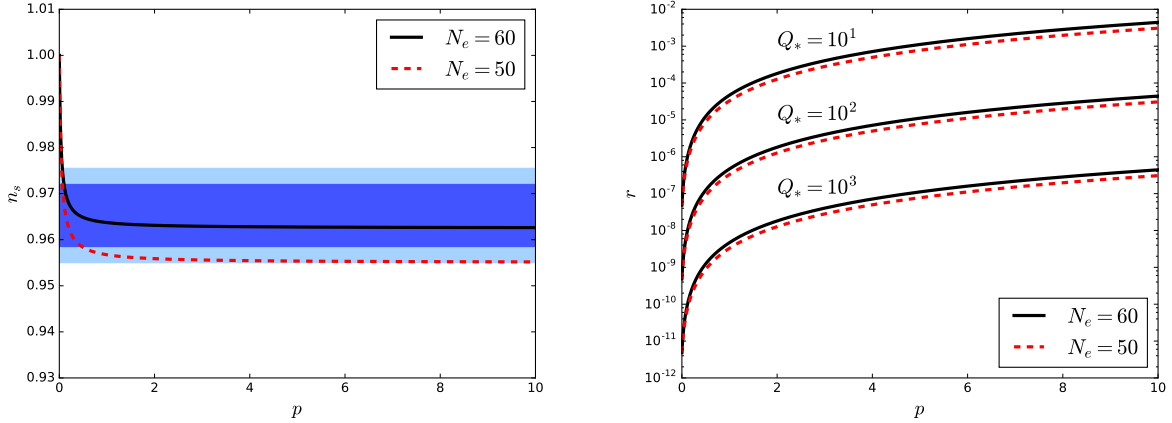


FIG. 3: Observational predictions of the DM scenario with a dissipative coefficient of the form $\Upsilon \propto \phi^p$ described by a quartic potential for 50 (red dashed-line) and 60 (black solid-line) e-folds of inflation. The plot on the left shows the spectral index n_s as a function the parameter p , while the plot on the right shows the tensor-to-scalar ratio r as a function of p for three values of Q_* . The blue contours correspond to the 68% and 95% C.L. results from Planck 2018 TT,TE,EE+lowE+lensing data [1]. Note that n_s lies well inside the Planck contours at $p \geq 1$. On the other hand, r gets enhanced for larger p .

C. $\Upsilon = C_T T$ and $V(\phi) = \lambda\phi^4$

We can ensure a linear dissipative coefficient by considering a DM model with a constant no. of bosonic (fermionic) light degrees of freedom N_M during the evolution of the inflaton field. This is equivalent at having in the continuum limit a density of states that goes inversely with T , such that $N_M \simeq 2Tn(\phi)/g$. Using this, and performing

the sums in Eqs (22) and (26) in the continuum limit, we then have:

$$\Upsilon^S(\phi, T) = N_M C_T^S T, \quad C_T^S = \frac{2g^2 \sqrt{f}}{\pi} \ln \left(\frac{2}{\sqrt{f}} \right) \left[1 - \frac{f}{4} \text{ArcTan} \left(\frac{4}{f} \right) \right], \quad (48)$$

$$\Upsilon^F(\phi, T) = N_M C_T^F T, \quad (49)$$

where C_T^F was given in Eq. (26). The observational predictions will depend now on the couplings g and h , and the parameter N_M .

To ensure the consistency of the analysis, namely the computation of the dissipation coefficients, we must verify that both the bosonic and fermionic fields coupled to the inflaton are kept close to thermal equilibrium, thus requiring $\bar{\Gamma}_{\chi_i}, \bar{\Gamma}_{\psi_{\chi_i}} > H$. Since both average decay widths are proportional to the temperature and smaller than the latter for perturbative couplings, if the latter conditions are satisfied we also ensure that $T \gtrsim H$ throughout inflation, thus validating the flat space approximation used in computing the dissipative coefficients. For the quartic potential, the ratio T/H increases during inflation, being proportional to the dissipative ratio Q , such that we only have to ensure that these conditions are met at the moment of horizon-crossing for the relevant CMB scales.

We note that in another realization of warm inflation with light fields, the Warm Little Inflaton (WLI) scenario, one has to impose additional constraints on the temperature during warm inflation so as to keep it above the mass of the fermionic fields coupled to the inflaton and below the underlying symmetry breaking scale. Such conditions are absent in the DM model, which is therefore less constraining. This comes of course at the expense of considering a whole tower of fields coupled to the inflaton, although only a small number of fields contributes effectively to dissipation at any given time.

In our analysis, we have also included the one loop thermal effective potential, computed in Appendix C. The bosonic contribution is given in Eq. (C9), while the fermionic one is given in Eq. (C6). From those expressions, one can see that the contribution of the finite temperature corrections to the derivatives of the effective potential is absent in this model. We have kept anyway the subdominant thermal contribution to the potential, setting the renormalisation scale such as $\mu = \exp(c_f/2)T_*$, where T_* is the temperature at horizon crossing. It is worth mentioning that by choosing the renormalisation scale through another prescription, for instance $\mu = \exp(c_b/2)T_*$, the outcome is not substantial altered.

With a T -dependent dissipative coefficient, the amplitude of the primordial spectrum Eq. (8) includes the growing mode function $G(Q_*)$. For a chaotic quartic potential, it was obtained numerically in [34, 36]:

$$G(Q_*) \simeq 1 + 0.0185 Q_*^{2.315} + 0.335 Q_*^{1.364}, \quad (50)$$

which alongside the dynamical evolution equations described in the previous section allows us to compute the inflationary observables for the chaotic model for 50 – 60 e-folds of inflation, provided that the adiabatic conditions discussed above are satisfied. The inflaton self-coupling λ is fixed as before by using the observed amplitude of curvature perturbations, $\Delta_R^2 \simeq 2.2 \times 10^{-9}$ [1], at 50-60 e-folds before the end of inflation.

For the quartic potential, we show in Fig. 4 the obtained values of the dissipative ratio at horizon-crossing, Q_* , as a function of the coupling g , for different numbers of light states and values of the Yukawa coupling h . We note in this figure that the adiabatic condition $\bar{\Gamma}_{\chi_i}/H \gtrsim 1$ for the fermions is only satisfied for $g \gtrsim 0.01$. We also find a lower bound $Q_* > 0.01$ for the smallest number of fields considered, and that the adiabatic conditions imply generically $h \gtrsim 1$. We only show in this figure values of $Q_* < 11$ since, as we will see below, larger values are incompatible with observational data [1].

This behaviour had been reported before for the WLI model [34–36]. For instance, the analysis in [36] presented a rather similar lower bound on $g \gtrsim 0.01$ and $Q_* \gtrsim 0.001$ due to the adiabatic condition, while the conditions on the temperature limit this coupling from above, $g \lesssim 0.1$ such that $Q_* \lesssim 0.25$. However, the DM model yields a substantially wider consistent parametric range, particularly for the value of the dissipative ratio at horizon-crossing.

Note that for larger values of the number of light states N_M and of the Yukawa coupling h we obtain a much narrower range of consistent values for the coupling g , and the lower bound on Q_* also increases with N_M and h . We thus find scenarios where inflation can start either in the weak or strong dissipation regimes, noting that Q grows during inflation such that for a wide range of parameters one reaches $Q > 1$ before the end of inflation, a necessary condition for radiation to dominate after the slow-roll regime with no further reheating (see eq. (6)).

In Figs. 5, 6 and 7 we show the predictions for the scalar spectral index and tensor-to-scalar ratio in the allowed parametric ranges, for different values of h and N_M , exhibiting a remarkable consistency with the Planck legacy

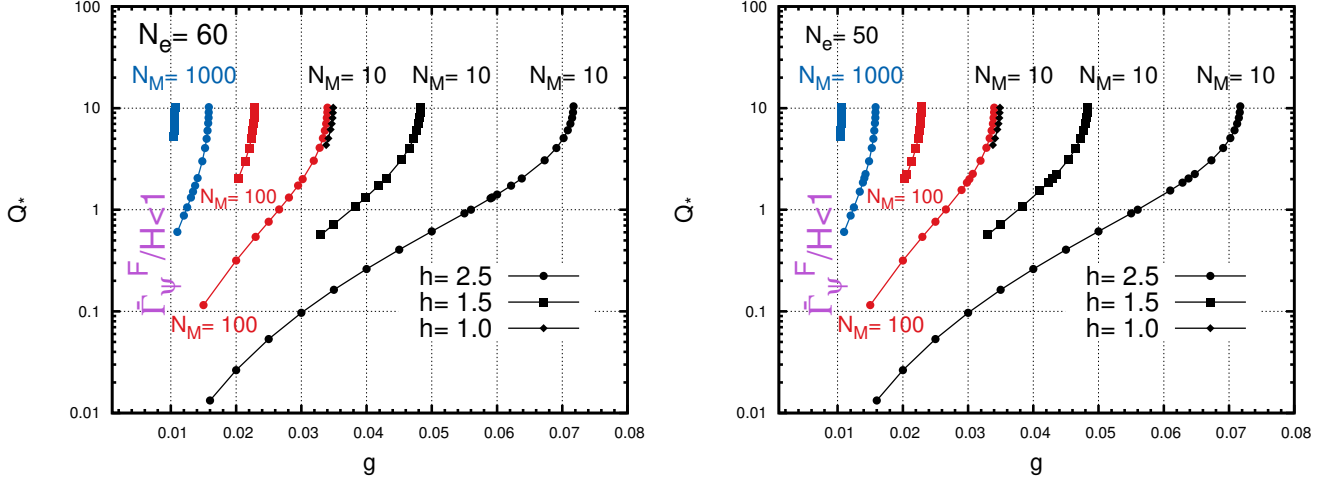


FIG. 4: Dissipative ratio at horizon-crossing, Q_* , as a function of the coupling g , within a DM scenario with a dissipative coefficient of the form $\Upsilon \propto T$ described by a quartic potential and $N_e = 50$ (right) or $N_e = 60$ (left) e-folds of inflation. We consider different numbers of light states $N_M = 10, 100$ and 1000 , shown by the black, red and blue curves, respectively; as well as different values of the Yukawa coupling $h = 2.5$ (circles), $h = 1.5$ (rectangles), and $h = 1.0$ (diamonds). Note that the adiabatic conditions $\bar{\Gamma}_{\chi_i}/H \gtrsim 1$ is not satisfied for small values of g .

results. This is particularly relevant given that the quartic potential is already excluded by such survey within the CI paradigm.

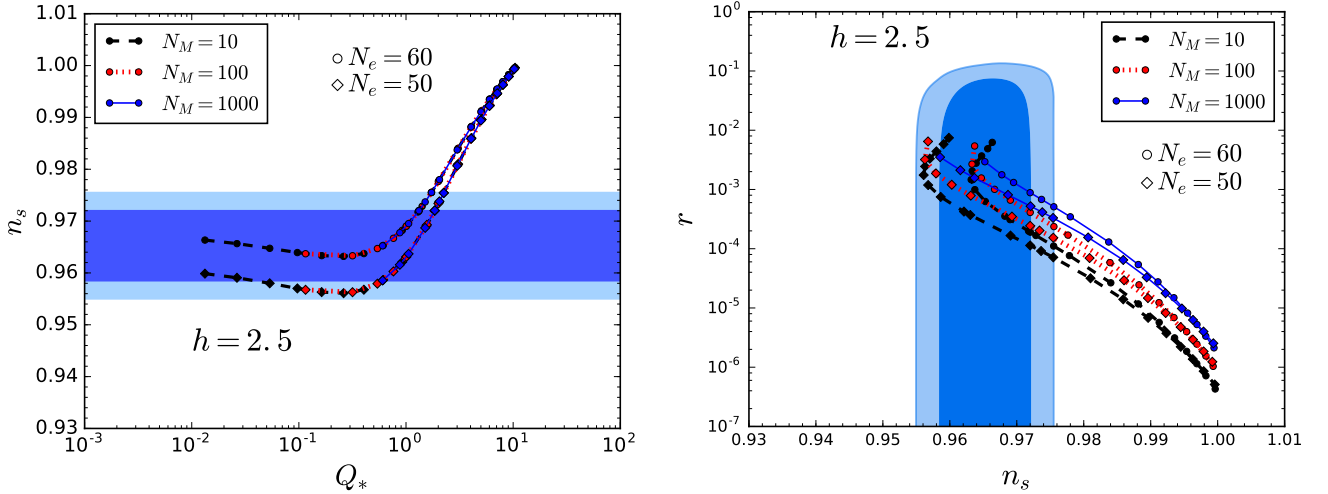


FIG. 5: Observational predictions of the DM scenario with a dissipative coefficient of the form $\Upsilon \propto T$ described by a quartic potential for 50 (diamonds) and 60 (circles) e-folds of inflation, a Yukawa coupling $h = 2.5$, and three different values of the number of light states $N_M = 10, 100, 1000$ in colours (line-style) black (dashed-line), red (pointed-line), and blue (solid-line) respectively. The plot on the left shows the spectral index n_s as a function of the dissipative ratio at horizon-crossing, Q_* , while the plot on the right shows the allowed trajectories in the (n_s, r) plane. The blue contours correspond to the 68% and 95% C.L. results from Planck 2018 TT,TE,EE+lowE+lensing data [1].

We find agreement with the Planck legacy data in the parametric ranges yielding values of $Q_* \lesssim 1$, which is easier to achieve for larger values of the Yukawa coupling, as well as a smaller number of light fields. For instance, we cannot find consistency with the Planck data for $h = 1$, since smaller values of the Yukawa coupling lead to stronger

dissipation at horizon-crossing. These results were expected, since the growing mode in the spectrum makes it more blue-tilted with increasing Q_* .

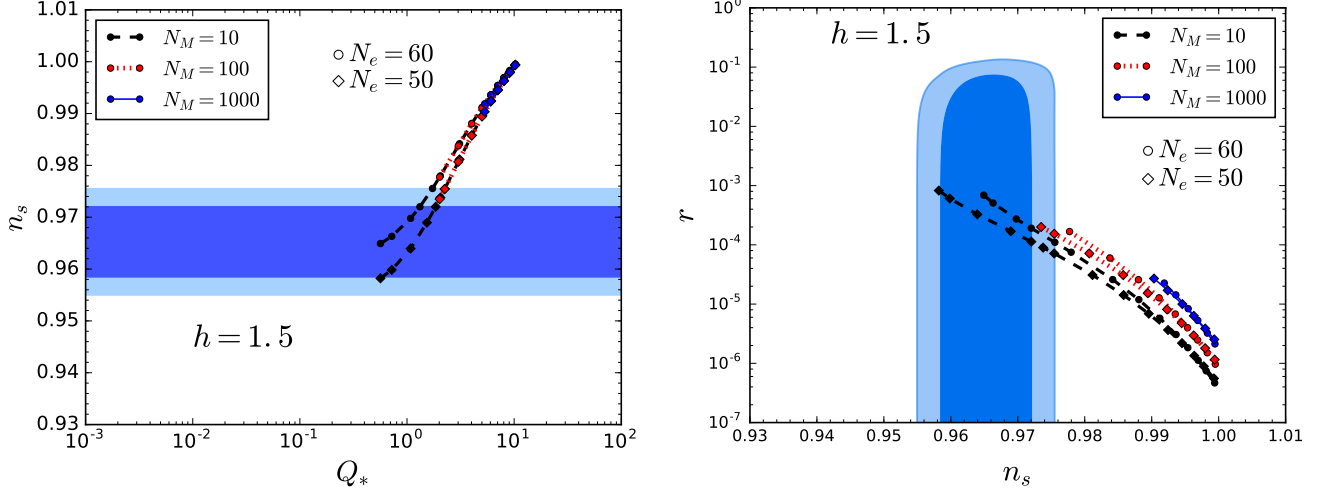


FIG. 6: Observational predictions of the DM scenario with a dissipative coefficient of the form $\Upsilon \propto T$ described by a quartic potential for 50 (diamonds) and 60 (circles) e-folds of inflation, a Yukawa coupling $h = 1.5$, and three different values of the number of light states $N_M = 10, 100, 1000$ in colours (line-style) black (dashed-line), red (pointed-line), and blue (solid-line) respectively. The plot on the left shows the spectral index n_s as a function of the dissipative ratio at horizon-crossing, Q_* , while the plot on the right shows the allowed trajectories in the (n_s, r) plane. The blue contours correspond to the 68% and 95% C.L. results from Planck 2018 TT,TE,EE+lowE+lensing data [1].

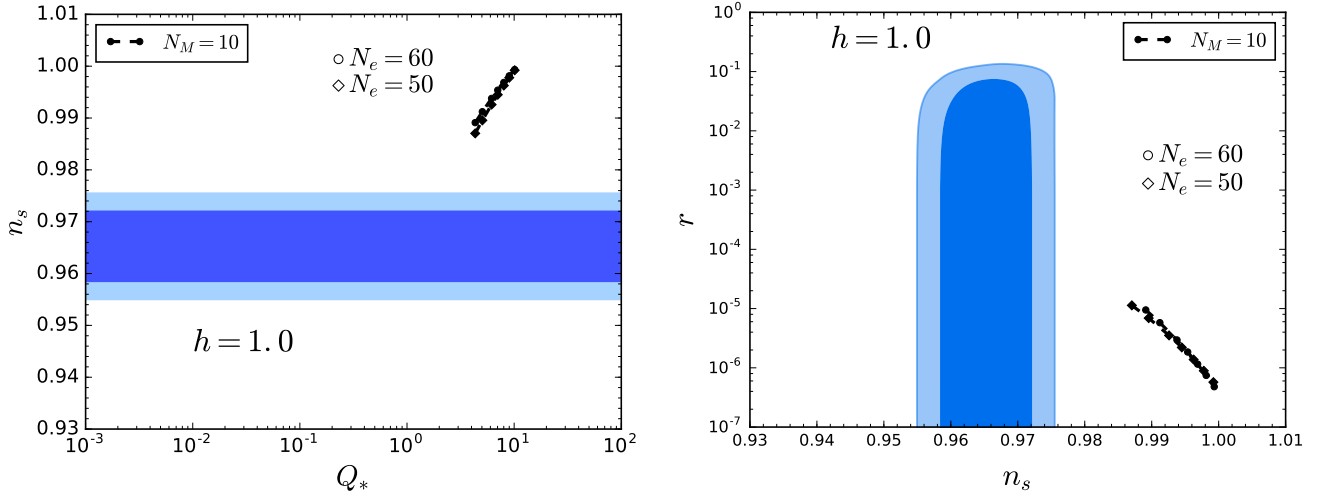


FIG. 7: Observational predictions of the DM scenario with a dissipative coefficient of the form $\Upsilon \propto T$ described by a quartic potential for 50 (diamonds) and 60 (circles) e-folds of inflation, a Yukawa coupling $h = 2.5$, and three different values of the number of light states $N_M = 10, 100, 1000$ in colours (line-style) black (dashed-line), red (pointed-line), and blue (solid-line) respectively. The plot on the left shows the spectral index n_s as a function of the dissipative ratio at horizon-crossing, Q_* , while the plot on the right shows the allowed trajectories in the (n_s, r) plane. The blue contours correspond to the 68% and 95% C.L. results from Planck 2018 TT,TE,EE+lowE+lensing data [1].

In general, the tensor-to-scalar ratio lies in the range $10^{-4} \lesssim r \lesssim 10^{-2}$, for values of the scalar spectral index within the Planck window, being more suppressed for larger values of Q_* (smaller Yukawa couplings and larger

number of light fields N_M). This feature is quite generic, since a similar notable trend was obtained in [25, 32–36]. Additionally, this significant agreement with Planck data is easily achieved with only a small number of light states N_M .

The remarkable agreement between the DM quartic model and the Planck data is similar to what had been obtained for the WLI quartic model [34–36], where the dissipation coefficient is also proportional to the temperature. However, while the more constraining conditions of the WLI model limit the consistent parametric range to $Q_* \lesssim 0.25$, we have found consistent DM scenarios with up to $Q_* \sim 1.85 - 2.25$ ($1.32 - 1.73$) for 50 (60) e-folds of inflation within the contours correspond to the 68% and 95% C.L. results from Planck 2018 TT,TE,EE+lowE+lensing data [1]. In addition, this noteworthy consistency with Planck, which also leads to a finite range for the tensor-to-scalar ratio, might be relevant in the search of primordial gravitational waves via B-mode polarization experiments in the near future (see e.g. [46]).

Finally we evaluate both swampland criteria within the observational window provided by Planck [1], and we find that none is fully satisfied, with $10 \lesssim |\Delta\phi|/M_P \lesssim 21$ and $0.19 \lesssim M_P|V_{eff,\phi}|/V_{eff} \lesssim 0.4$ for 60 e-folds and $7 \lesssim |\Delta\phi|/M_P \lesssim 19$ and $0.2 \lesssim M_P|V_{eff,\phi}|/V_{eff} \lesssim 0.55$ for 50 e-folds. Nevertheless, the scenarios for which $Q_* \sim 1$ are closer to satisfying these criteria than the corresponding cold inflation scenarios, and could in fact satisfied them if they were somewhat relaxed. For instance, it has been discussed in [18, 47] that c could be as small as $\mathcal{O}(10^{-11})$, and it does not object perceiving de Sitter vacua in string landscapes. Moreover, there has been some controversy on the precise values for Δ and c , given that they depend on the specific string model being considered [48]. Basically, if Δ were increased by an order of magnitude, and c were decreasing by around $c \sim 0.4$ both swampland criteria could be easily fulfilled for $Q_* \gtrsim 1$.

VII. CONCLUSIONS

Warm inflation offers a unique description of the early universe cosmology. It is a picture of an inflationary dynamics in which the state of the universe during inflation is not the vacuum state, but rather an excited statistical thermal state. It introduces dissipation into the inflationary dynamics which can be well explained by first principles of a quantum multi-field theory. Besides, this approach has several attractive features. For instance, the additional friction may ease the required flatness of the inflaton potential alleviating the swampland criterion. Also, even if radiation is subdominant during inflation, it may smoothly become the leading component if $Q \gtrsim 1$ at the end of inflation ($\epsilon_{eff} \sim 1 + Q$), with no need for a separate reheating period^e. It also may explain the nature of the classical inhomogeneities observed in the CMB, since for WI the fluctuations of the inflaton are thermally induced; hence there is no need to explain the troublesome quantum-to-classical transition problem of CI, due to the purely quantum origin of the CI density perturbations. Furthermore, it was shown early on [10], in fact in studying the DM model, that the presence of dissipation and temperature lowers the energy scale of inflation in monomial models such as $\lambda\phi^4$ in comparison to the cold inflation results. This was further developed in [25, 32, 33], with the prediction for a low tensor-to-scalar ratio, which subsequently has been shown to be consistent with data [1]. These observational predictions for warm inflation were most recently tested in [49] with good agreement with Planck data.

The ultimate goal in building inflation models is to have them consistent with observational data and also be theoretically consistent. The latter has many levels of criteria. In [50], based on general properties of quantum field theory, conditions were listed for theoretically consistent warm inflation models. In fact a subset of those conditions would also apply to cold inflation models. Amongst the criteria stated were that $\phi < M_P$ and $m_\phi > H$. The former was based on general consideration of low-energy effective field theory. The latter is so that inflaton particles remained sub-Hubble scale and is basically the η -condition. However in contrast to cold inflation, where this condition is something one needs to model build around, in warm inflation due to the presence of dissipation this condition can hold and inflation can still occur. In Sect. (VI) explicit examples of warm inflation models are

^e Although we have not specified how the fields in the tower interact with the Standard Model degrees of freedom, which is a model-dependent issue, these will generically be excited at the end of inflation once the Hubble rate is sufficiently reduced, given the large values of the temperature $T \sim 10^{12-13}$ GeV when radiation becomes dominant. In fact, it would require extremely small couplings between the tower and Standard Model states to prevent the latter's thermalization before nucleosynthesis at $T \sim 10$ MeV (see also [55]).

shown where $\phi < M_p$ and $m_\phi > H$, such as in Fig. 2. These are very stringent requirements for an inflation model and no cold inflation models can achieve both of them. The swampland criteria are basically contained within these criteria. In this paper we have found warm inflation models which are consistent with all the criteria stated in [50] and so in turn are also consistent with the swampland criteria. Supergravity corrections will also, in principle, be under control in such supersymmetric scenarios with sub-Planckian field values, although scenarios with $\phi > M_p$ may also be viable in supergravity with non-canonical choices of the Kähler potential (see e.g. [56]).

The distributed mass model of warm inflation studied in this paper develops the original studies to include different field- and temperature-dependent mass distributions and makes accurate comparisons with the most recent Planck data. This model fits naturally with the string landscape picture, with different mass distributions in the model ultimately arising from different possible string vacua. We found for example for the $\Upsilon \sim \phi$ model with $\lambda\phi^4$ potential, full consistency with the conditions in [50] and therefore also consistency with the swampland criteria. As the $\lambda\phi^4$ potential in many respects is the benchmark inflation model, this consistency is a very promising result both for warm inflation and the DM implementation of landscape phenomenology. We also examined other types of mass distribution functions with varied level of success. Overall this shows these models have a robust range of possibility. These results encourage further understanding of DM models in particular in the context of string theory building on the ideas developed in [9].

If one accepts string theory and thus its landscape property as the fundamental description of nature, then it forces a rethink on the relevance of simplicity for an inflation model. Any point on the landscape that leads to a consistent effective field theory model which agrees with observational data is just as good as any other. We refrain from calling this what it is, but that is the path this line of reasoning evidently takes us down. From the present vantage point with so many possible vacua, it would seem easy to argue that any working low-energy effective model leading to observationally consistent inflation can be found somewhere in this vast landscape. Alternatively said, it is hard to argue that simple low-energy models are the only ones that could arise from this landscape. It is possible that future work in string theory will reduce the number of viable vacua, but given the starting point, that seems a difficult task. From another viewpoint, the loss of predictivity that comes with the landscape property, equally encourages objective thinking whether string theory is the best way to approach the ultraviolet completion problem.

In a certain way of looking at it, DM models are rather complicated in that there are many fields and several conditions to examine and calculate. However from another perspective they are also simple in that they are renormalizable models, with canonical kinetic terms and require no assumptions of the coupling of gravity. Due to these properties, they are very amenable to be part of an extension of the Standard Model. All said, from the landscape perspective, DM models are viable models that are interesting to further explore.

Acknowledgments

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Appendix A: Dissipative coefficient in the high temperature regime: pole approximation

The leading contribution to the dissipative coefficient from the scalar χ_i fields has the following form [39, 40]:

$$\Upsilon = \sum_{i=1}^{N_M} \frac{4}{T} \left(\frac{g^2}{2} \right)^2 (\phi - M_i)^2 \int \frac{d^4 p}{(2\pi)^4} \rho_\chi^2 n_B (1 + n_B), \quad (\text{A1})$$

where $n_B = [e^{p_0/T} - 1]^{-1}$ is the Bose-Einstein distribution and ρ_χ is the spectral function for the χ_i field. The spectral function for the χ_i field entering in eq.(A1) corresponds to the fully dressed propagator, including the effect of its finite decay width into σ_i and ψ_{σ_i} particles:

$$\rho_\chi(p, p_0) = \frac{4\omega_p \Gamma_{\chi_i}}{(p_0^2 - \omega_p^2)^2 + 4\omega_p^2 \Gamma_{\chi_i}^2}, \quad (\text{A2})$$

where $\omega_p = \sqrt{\tilde{m}_{\chi_i}^2 + p^2}$ and $\tilde{m}_{\chi_i}^2 = m_{\chi_i}^2 + g^2/12 + h^2 T^2/8$. The leading contribution to the decay width of the χ_i fields is then the two-body decay $\chi_i \rightarrow \sigma_i \sigma_i$ and $\chi_i \rightarrow \bar{\psi}_{\sigma_i} \psi_{\sigma_i}$, where at finite temperature we include contributions from both decays and inverse decays, as well as thermal scattering off particles in the radiation bath. The first decay process has been computed in [39, 40] from the imaginary part of the χ_i self energy at one-loop order, yielding:

$$\Gamma_{\chi_i}(\chi_i \rightarrow \sigma_i \sigma_i) = \frac{h^2}{32\pi} \frac{m_{\chi_i}^2}{\omega_p} F_T^{\sigma_i \sigma_i}(p, p_0), \quad (\text{A3})$$

where

$$\begin{aligned} F_T^{\sigma_i \sigma_i}(p, p_0) &= \left[\frac{\omega_+^{\sigma_i} - \omega_-^{\sigma_i}}{p} + \frac{T}{p} \ln \left(\frac{1 - e^{-\frac{\omega_+^{\sigma_i}}{T}}}{1 - e^{-\frac{\omega_-^{\sigma_i}}{T}}} \frac{1 - e^{-\frac{p_0 - \omega_-^{\sigma_i}}{T}}}{1 - e^{-\frac{p_0 - \omega_+^{\sigma_i}}{T}}} \right) \right] \theta(p_0^2 - p^2 - 4m_{\sigma_i}^2) \\ &+ \left[\frac{T}{p} \ln \left(\frac{1 - e^{-\frac{\omega_+^{\sigma_i}}{T}}}{1 - e^{-\frac{\omega_-^{\sigma_i}}{T}}} \frac{1 - e^{-\frac{p_0 + \omega_-^{\sigma_i}}{T}}}{1 - e^{-\frac{p_0 + \omega_+^{\sigma_i}}{T}}} \right) \right] \theta(-p_0^2 + p^2), \end{aligned} \quad (\text{A4})$$

with m_{σ_i} denoting the scalar σ_i mass, $\theta(x)$ the Heaviside function and

$$\omega_{\pm}^{\sigma_i} = \sqrt{(k_{\pm}^{\sigma_i})^2 + m_{\sigma_i}^2}, \quad k_{\pm}^{\sigma_i} = \frac{1}{2} \left| p \pm p_0 \left(1 - \frac{4m_{\sigma_i}^2}{p_0^2 - p^2} \right)^{1/2} \right|. \quad (\text{A5})$$

On the other hand, the leading contribution to the decay width of the χ_i fields into two fermions $\bar{\psi}_{\sigma_i} \psi_{\sigma_i}$, has been computed in [39] from the imaginary part of the χ_i self energy at one-loop order, yielding:

$$\Gamma_{\chi_i}(\chi_i \rightarrow \bar{\psi}_{\sigma_i} \psi_{\sigma_i}) = \frac{h^2}{32\pi\omega_p} F_T^{\bar{\psi}_{\sigma_i} \psi_{\sigma_i}}(p, p_0), \quad (\text{A6})$$

where

$$\begin{aligned} F_T^{\bar{\psi}_{\sigma_i} \psi_{\sigma_i}}(p, p_0) &= (p_0^2 - p^2) \left[1 - \frac{4m_{\psi_{\sigma_i}}^2}{p_0^2 - p^2} \right] \left[\frac{\omega_+^{\psi_{\sigma_i}} - \omega_-^{\psi_{\sigma_i}}}{p} + \frac{T}{p} \ln \left(\frac{1 + e^{-\frac{\omega_+^{\psi_{\sigma_i}}}{T}}}{1 + e^{-\frac{\omega_-^{\psi_{\sigma_i}}}{T}}} \frac{1 + e^{-\frac{p_0 - \omega_-^{\psi_{\sigma_i}}}{T}}}{1 + e^{-\frac{p_0 - \omega_+^{\psi_{\sigma_i}}}{T}}} \right) \right] \theta(p_0^2 - p^2 - 4m_{\psi_{\sigma_i}}^2) \\ &+ (p_0^2 - p^2) \left[1 - \frac{4m_{\psi_{\sigma_i}}^2}{p_0^2 - p^2} \right] \frac{T}{p} \ln \left(\frac{1 + e^{-\frac{\omega_+^{\psi_{\sigma_i}}}{T}}}{1 + e^{-\frac{\omega_-^{\psi_{\sigma_i}}}{T}}} \frac{1 + e^{-\frac{p_0 + \omega_-^{\psi_{\sigma_i}}}{T}}}{1 + e^{-\frac{p_0 + \omega_+^{\psi_{\sigma_i}}}{T}}} \right) \theta(-p_0^2 + p^2), \end{aligned} \quad (\text{A7})$$

then $m_{\psi_{\sigma_i}}$ is the fermionic ψ_{σ_i} mass and

$$\omega_{\pm}^{\psi_{\sigma_i}} = \sqrt{(k_{\pm}^{\psi_{\sigma_i}})^2 + m_{\psi_{\sigma_i}}^2}, \quad k_{\pm}^{\psi_{\sigma_i}} = \frac{1}{2} \left| p \pm p_0 \left(1 - \frac{4m_{\psi_{\sigma_i}}^2}{p_0^2 - p^2} \right)^{1/2} \right|. \quad (\text{A8})$$

An analysis of Eq.(A1) indicates that the behaviour of the dissipation coefficient in the different temperature and interaction regimes is determined by an interplay between the spectral function and the thermal occupation numbers. When the temperature is large, the occupation numbers are also large, so in this regime, the poles of the spectral functions will dominate the integral and this will be referred to as the pole approximation. As such, the pole approximation works well in the high- T region. The dissipation coefficient about its pole at $p_0 \simeq \omega_p$ is:

$$\Upsilon = \sum_{i=1}^{N_M} \frac{2}{T} \left(\frac{g^2}{2} \right)^2 (\phi - M_i)^2 \int \frac{d^3 p}{(2\pi)^3} \frac{n_B(1 + n_B)}{\Gamma_{\chi_i} \omega_p^2}. \quad (\text{A9})$$

For on-shell χ_i modes, decays into light scalars and fermions are equally probable, with partial decay widths given by Eqs. (A3) and (A6). By treating the scalar σ_i and the fermion ψ_{σ_i} as light fields, we evaluate both decay rates in the limit $m_{\sigma_i} \ll T$ ($m_{\psi_{\sigma_i}} \ll T$), so $\omega_{\pm}^{\sigma_i}/T \simeq k_{\pm}^{\sigma_i}/T$ ($\omega_{\pm}^{\psi_{\sigma_i}}/T \simeq k_{\pm}^{\psi_{\sigma_i}}/T$); this estimation also yields that $k_{\pm}^{\sigma_i} = k_{\pm}^{\psi_{\sigma_i}} = |p_0 \pm p|/2$. Finally, in the limit $m_{\chi_i} \ll T$, we have that both partial decay widths are given by:

$$\Gamma_{\chi_i}(\chi_i \rightarrow \sigma_i \sigma_i) = \frac{h^2}{8\pi} \frac{m_{\chi_i}^2}{\tilde{m}_{\chi_i}} \frac{T}{\omega_p}, \quad (\text{A10})$$

$$\Gamma_{\chi_i}(\chi_i \rightarrow \bar{\psi}_{\sigma_i} \psi_{\sigma_i}) = \frac{h^2}{128\pi} \frac{\tilde{m}_{\chi_i}^3}{\omega_p T}, \quad (\text{A11})$$

yielding a total decay rate:

$$\Gamma_{\chi_i} = \frac{h^2}{128\pi} \frac{T \tilde{m}_{\chi_i}}{\omega_{\chi_i}(p)} \left[16 \frac{m_{\chi_i}^2}{\tilde{m}_{\chi_i}^2} + \frac{\tilde{m}_{\chi_i}^2}{T^2} \right]. \quad (\text{A12})$$

Finally we compute the integral in Eq. (A9), by following the same procedure as in [8], obtaining the total scalar dissipative coefficient:

$$\Upsilon^S(\phi, T) = \sum_{i=1}^{N_M} \frac{32g^4}{\pi h^2 \left[16 \frac{m_{\chi_i}^2}{\tilde{m}_{\chi_i}^2} + \frac{\tilde{m}_{\chi_i}^2}{T^2} \right]} \ln \left(\frac{2T}{\tilde{m}_{\chi_i}} \right) \frac{(\phi - M_i)^2}{\tilde{m}_{\chi_i}}. \quad (\text{A13})$$

Appendix B: Scalar decay width average

The thermal average of the scalar decay width is given by:

$$\bar{\Gamma}_{\chi_i} = \frac{1}{n_B} \int \frac{d^3p}{(2\pi)^3} \Gamma_{\chi_i} f_B, \quad (\text{B1})$$

where $f_B(\mathbf{p}/T)$ is the Bose-Einstein distribution and $n_B(\tilde{m}_{\chi_i}/T)$ the associated number density. Hence we have:

$$\bar{\Gamma}_{\chi_i} = \frac{h^2 T \tilde{m}_{\chi_i}}{128\pi} \left[16 \frac{m_{\chi_i}^2}{\tilde{m}_{\chi_i}^2} + \frac{\tilde{m}_{\chi_i}^2}{T^2} \right] I(\tilde{m}_{\chi_i}/T), \quad I(\tilde{m}_{\chi_i}/T) = \frac{\int \frac{d^3p}{(2\pi)^3} \frac{f_B}{\omega_{\chi_i}(p)}}{\int \frac{d^3p}{(2\pi)^3} f_B}. \quad (\text{B2})$$

The integral factor $I(\tilde{m}_{\chi_i}/T)$ can be obtained numerically and is well approximated by:

$$I(\tilde{m}_{\chi_i}/T) = \frac{1}{T} \frac{\int \frac{dx x^2}{\sqrt{x^2+a^2}} (e^{\sqrt{x^2+a^2}} - 1)^{-1}}{\int dx x^2 (e^{\sqrt{x^2+a^2}} - 1)^{-1}} \simeq \frac{1}{T} \frac{0.68}{(1 + 0.77a)}, \quad (\text{B3})$$

where we have defined $x = p/T$ and $a = \tilde{m}_{\chi_i}/T$. Hence the thermal average of the scalar decay width is:

$$\bar{\Gamma}_{\chi_i} \simeq \frac{h^2 \tilde{m}_{\chi_i}}{128\pi} \left[16 \frac{m_{\chi_i}^2}{\tilde{m}_{\chi_i}^2} + \frac{\tilde{m}_{\chi_i}^2}{T^2} \right] \frac{0.68}{(1 + 0.77f^{1/2})}. \quad (\text{B4})$$

Moreover, in the limit $m_{\chi_i}/T \lesssim 1$ we have that the effective finite thermal mass \tilde{m}_{χ_i} can be fittingly taken as $\tilde{m}_{\chi_i} \simeq f^{1/2}T$, where $f = f(g, h) = g^2/12 + h^2/8$. Therefore the average decay width becomes:

$$\bar{\Gamma}_{\chi_i} \simeq \frac{h^2 f^{1/2}}{128\pi} \left[\frac{16}{f} \frac{m_{\chi_i}^2}{T^2} + f \right] \frac{0.68}{(1 + 0.77f^{1/2})} T. \quad (\text{B5})$$

Appendix C: Effective potential at finite temperature

We start by considering the contribution of the light fermions in the tower to the finite temperature effective potential, which is given by [51, 53, 54]:

$$V_T^{\psi_{\chi_i}} \simeq \frac{1}{2} \sum_i \left\{ -\frac{7\pi^2}{180} T^4 + \frac{\tilde{m}_{\psi_{\chi_i}}^2 T^2}{12} + \frac{\tilde{m}_{\psi_{\chi_i}}^4}{16\pi^2} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right\}, \quad (\text{C1})$$

where μ is the $\overline{\text{MS}}$ renormalization scale, $c_f = 2.635$, and the effective thermal fermion masses $\tilde{m}_{\psi_i}^2 = g^2(\phi - M_i)^2 + h^2 T^2/8$. Note that the overall factor $1/2$ in front of the sum is related to the Majorana nature of the fermions in the SUSY model. In the continuum limit we may write this in the form:

$$\begin{aligned} V_T^{\psi_i} &\simeq \frac{T^4}{2} \sum_i \left\{ -\frac{7\pi^2}{180} + \left[\frac{g^2(\phi - M_i)^2}{T^2} + \frac{h^2}{8} \right] \frac{1}{12} \right. \\ &\quad \left. + \left[\frac{g^4(\phi - M_i)^4}{T^4} + \frac{h^2 g^2(\phi - M_i)^2}{4 T^2} + \frac{h^4}{64} \right] \frac{1}{16\pi^2} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right\} \\ &\simeq \frac{T^4}{2} \left\{ \left[-\frac{7\pi^2}{180} + \frac{h^2}{96} + \frac{h^4}{1024\pi^2} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right] \int_{M_-}^{M_+} dM n(M) \right. \\ &\quad \left. + \frac{g^2}{T^2} \left[\frac{1}{12} + \frac{h^2}{64\pi^2} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right] \int_{M_-}^{M_+} dM n(M) (\phi - M)^2 \right. \\ &\quad \left. + \frac{1}{16\pi^2} \frac{g^4}{T^4} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \int_{M_-}^{M_+} dM n(M) (\phi - M)^4 \right\} \\ &\simeq \frac{T^4}{2} \left\{ \left[-\frac{7\pi^2}{180} + \frac{h^2}{96} + \frac{h^4}{1024\pi^2} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right] \left[\frac{2T}{g} n(\phi) + \dots \right] \right. \\ &\quad \left. + \frac{g^2}{T^2} \left[\frac{1}{12} + \frac{h^2}{64\pi^2} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right] \left[\frac{2}{3} \frac{T^3}{g^3} n(\phi) + \dots \right] \right. \\ &\quad \left. + \frac{1}{16\pi^2} \frac{g^4}{T^4} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \left[\frac{2}{5} \frac{T^5}{g^5} n(\phi) + \dots \right] \right\} \\ &\simeq \frac{T^5}{2g} n(\phi) \left\{ -\frac{7\pi^2}{90} + \frac{1}{18} + \frac{h^2}{48} + \frac{1}{16\pi^2} \left(\frac{2}{5} + \frac{h^2}{6} + \frac{h^4}{32} \right) \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right\} + \dots, \quad (\text{C2}) \end{aligned}$$

where the integrals are in general

$$\begin{aligned} \int_{M_-}^{M_+} dM n(M) (\phi - M)^n &= \int_{-T/g}^{T/g} dx \left[n(\phi) - x n'(\phi) + \frac{x^2}{2!} n''(\phi) + \dots \right] x^n \\ &= n(\phi) \frac{x^{n+1}}{n+1} - n'(\phi) \frac{x^{n+2}}{n+2} + n''(\phi) \frac{x^{n+3}}{n+3} + \dots \Big|_{-T/g}^{T/g} \\ &= \begin{cases} \frac{2}{n+1} \left(\frac{T}{g} \right)^{n+1} n(\phi) + \dots & , n = 0, 2, 4, \dots \\ -\frac{2}{n+2} \left(\frac{T}{g} \right)^{n+2} n'(\phi) + \dots & , n = 1, 3, 5, \dots \end{cases} \quad (\text{C3}) \end{aligned}$$

We have considered the local approximation, i.e. that only states in the vicinity of $M = \phi(t)$ are light at any given time. Note that to obtain the derivatives of this potential correction with respect to the field, one needs to take into account that both the integrand and the integration limits are ϕ -dependent, the end result corresponding to

differentiating eq. (C2), such that:

$$V_{T,\phi}^{\psi_i} \simeq \frac{T^5}{2g} n'(\phi) \left\{ -\frac{7\pi^2}{90} + \frac{1}{18} + \frac{h^2}{48} + \frac{1}{16\pi^2} \left(\frac{2}{5} + \frac{h^2}{6} + \frac{h^4}{32} \right) \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right\} + \dots, \quad (C4)$$

$$V_{T,\phi\phi}^{\psi_i} \simeq \frac{T^5}{2g} n''(\phi) \left\{ -\frac{7\pi^2}{90} + \frac{1}{18} + \frac{h^2}{48} + \frac{1}{16\pi^2} \left(\frac{2}{5} + \frac{h^2}{6} + \frac{h^4}{32} \right) \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right\} + \dots. \quad (C5)$$

Nevertheless, in the local approximation the density of states could be taken only as temperature-dependent $n \sim T^{-1}$, and independent of the field; therefore $n'(\phi) = n''(\phi) = 0$. Thus we have $N_M \simeq 2Tn/g$, hence $n \simeq gN_M/(2T)$; but most importantly $V_{T,\phi}^{\psi_{\chi_i}} = V_{T,\phi\phi}^{\psi_{\chi_i}} \simeq 0$. Finally the finite temperature corrections to the effective potential due to the fermionic tower, in the local approximation at leading order, can be written as:

$$V_T^{\psi_i} \simeq \frac{N_M}{4} T^4 \left\{ -\frac{7\pi^2}{90} + \frac{1}{18} + \frac{h^2}{48} + \frac{1}{16\pi^2} \left(\frac{2}{5} + \frac{h^2}{6} + \frac{h^4}{32} \right) \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_f \right] \right\}. \quad (C6)$$

Next we examine the contribution of the light bosons in the tower to the finite temperature effective potential given by [51, 53, 54]:

$$V_T^{\chi_i} \simeq 2 \sum_i \left\{ -\frac{\pi^2}{90} T^4 + \frac{\tilde{m}_{\chi_i}^2 T^2}{24} - \frac{\tilde{m}_{\chi_i}^3 T}{12\pi} - \frac{\tilde{m}_{\chi_i}^4}{64\pi^2} \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_b \right] \right\}, \quad (C7)$$

where again μ denotes the $\overline{\text{MS}}$ renormalization scale, $c_b = 5.41$, and the effective thermal boson masses $\tilde{m}_{\chi_i}^2 = g^2(\phi - M_i)^2 + g^2 T^2/12 + h^2 T^2/8$. Note the overall factor 2 in front of the sum, which represents the fact that the χ_i 's scalars are complex fields. Following the same procedure as for the fermionic tower, we have for the bosonic sector:

$$V_T^{\chi_i} \simeq \frac{2T^5}{g} n(\phi) \left\{ -\frac{\pi^2}{45} + \frac{f^2(g, h)}{16\pi} \ln \left(\frac{\sqrt{f(g, h)}}{\sqrt{1 + f(g, h)} + 1} \right) + \frac{f(g, h)}{12} - \frac{(2 + 5f(g, h))}{48\pi} \sqrt{1 + f(g, h)} \right. \\ \left. - \frac{1}{16\pi^2} \left(\frac{1}{10} + \frac{f(g, h)}{3} + \frac{f^2(g, h)}{2} \right) \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_b \right] \right\} + \dots, \quad (C8)$$

where $f(g, h) = g^2/12 + h^2/8$. In the local approximation $n(\phi) \sim T^{-1}$, hence $n(\phi) \simeq gN_\chi/(2T)$, and the inflaton ϕ derivatives are zero: $V_{T,\phi}^{\chi_i} = V_{T,\phi\phi}^{\chi_i} \simeq 0$. Finally the finite temperature corrections to the effective potential due to the bosonic tower, in the local approximation at leading order, can be written in the form:

$$V_T^{\chi_i} \simeq N_M T^4 \left\{ -\frac{\pi^2}{45} + \frac{f^2(g, h)}{16\pi} \ln \left(\frac{\sqrt{f(g, h)}}{\sqrt{1 + f(g, h)} + 1} \right) + \frac{f(g, h)}{12} - \frac{(2 + 5f(g, h))}{48\pi} \sqrt{1 + f(g, h)} \right. \\ \left. - \frac{1}{16\pi^2} \left(\frac{1}{10} + \frac{f(g, h)}{3} + \frac{f^2(g, h)}{2} \right) \left[\ln \left(\frac{\mu^2}{T^2} \right) - c_b \right] \right\}. \quad (C9)$$

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